

# Production of doubly charged scalars from the decay of singly charged scalars in the Higgs Triplet Model

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(Dated: January 20, 2013)

## Abstract

The existence of doubly charged Higgs bosons ( $H^{\pm\pm}$ ) is a distinctive feature of the Higgs Triplet Model (HTM), in which neutrinos obtain tree-level masses from the vacuum expectation value of a neutral scalar in a triplet representation of  $SU(2)_L$ . We point out that a large branching ratio for the decay of a singly charged Higgs boson to a doubly charged Higgs boson via  $H^\pm \rightarrow H^{\pm\pm}W^*$  is possible in a sizeable parameter space of the HTM. From the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  the above decay mode would give rise to pair production of  $H^{\pm\pm}$ , with a cross section which can be comparable to that of the standard pair-production mechanism  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ . We suggest that the presence of a sizeable branching ratio for  $H^\pm \rightarrow H^{\pm\pm}W^*$  could significantly enhance the detection prospects of  $H^{\pm\pm}$  in the four-lepton channel. Moreover, the decays  $H^0 \rightarrow H^\pm W^*$  and  $A^0 \rightarrow H^\pm W^*$  from production of the neutral triplet scalars  $H^0$  and  $A^0$  would also provide an additional source of  $H^\pm$ , which can subsequently decay to  $H^{\pm\pm}$ .

PACS numbers: 14.80.Fd, 12.60.Fr, 14.60.Pq

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## I. INTRODUCTION

The established evidence that neutrinos oscillate and possess small masses [1] necessitates physics beyond the Standard Model (SM), which could manifest itself at the CERN Large Hadron Collider (LHC) and/or in low energy experiments which search for the lepton flavour violation [2]. Consequently, models of neutrino mass generation which can be probed at present and forthcoming experiments are of great phenomenological interest.

Neutrinos may obtain mass via the vacuum expectation value (vev) of a neutral Higgs boson in an isospin triplet representation [3–7]. A particularly simple implementation of this mechanism of neutrino mass generation is the “Higgs Triplet Model” (HTM) in which the SM Lagrangian is augmented solely by  $\Delta$  which is a  $SU(2)_L$  triplet of scalar particles with hypercharge  $Y = 2$  [3, 6, 7]. In the HTM, the Majorana neutrino mass matrix  $m_{\ell\ell'}$  ( $\ell, \ell' = e, \mu, \tau$ ) is given by the product of a triplet Yukawa coupling matrix  $h_{\ell\ell'}$  and a triplet vev ( $v_\Delta$ ). Consequently, the direct connection between  $h_{\ell\ell'}$  and  $m_{\ell\ell'}$  gives rise to phenomenological predictions for processes which depend on  $h_{\ell\ell'}$  because  $m_{\ell\ell'}$  has been restricted well by neutrino oscillation measurements [1, 8–12]. A distinctive signal of the HTM would be the observation of doubly charged Higgs bosons ( $H^{\pm\pm}$ ) whose mass ( $m_{H^{\pm\pm}}$ ) may be of the order of the electroweak scale. Such particles can be produced with sizeable rates at hadron colliders in the processes  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  [13–17] and  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  [13, 18, 19]. The first searches for  $H^{\pm\pm}$  at a hadron collider were carried out at the Fermilab Tevatron, assuming the production channel  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  and decay  $H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm$ . The mass limits  $m_{H^{\pm\pm}} > 110 \rightarrow 150$  GeV [20, 21] were derived, with the strongest limits being for  $\ell = e, \mu$  [20]. The branching ratios (BRs) for  $H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm$  depend on  $h_{\ell\ell'}$  and are predicted in the HTM in terms of the parameters of the neutrino mass matrix [19, 22, 23]. Detailed quantitative studies of  $\text{BR}(H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm)$  in the HTM have been performed in [24–27] with particular emphasis given to their sensitivity to the Majorana phases and the absolute neutrino mass i.e. parameters which cannot be probed in neutrino oscillation experiments. A study on the relation between  $\text{BR}(H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm)$  and the neutrinoless double beta decay can be seen in [28]. Simulations of the detection prospects of  $H^{\pm\pm}$  at the LHC with  $\sqrt{s} = 14$  TeV previously focussed on  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  only [29], but recent studies now include the mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  [27, 30, 31]. The first search for  $H^{\pm\pm}$  at the LHC with  $\sqrt{s} = 7$  TeV [32] has recently been performed for both production mechanisms  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  and  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$ , for the decay channels  $H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm$  and  $H^\pm \rightarrow \ell^\pm\nu_{\ell'}$ .

In phenomenological studies of the HTM, for simplicity it is sometimes assumed that  $H^{\pm\pm}$  and  $H^\pm$  are degenerate, with a mass  $M$  which arises from a bilinear term  $M^2\text{Tr}(\Delta^\dagger\Delta)$  in the scalar potential. In this scenario the only possible decay channels for  $H^{\pm\pm}$  are  $H^{\pm\pm} \rightarrow \ell^\pm\ell'^\pm$  and  $H^{\pm\pm} \rightarrow W^\pm W^\pm$ , and the branching ratios are determined by the magnitude of  $v_\Delta$ . However, quartic terms in the scalar potential break the degeneracy of  $H^{\pm\pm}$  and  $H^\pm$ , and induce a mass splitting  $\Delta M \equiv m_{H^{\pm\pm}} - m_{H^\pm}$ , which can be of either sign. If  $m_{H^{\pm\pm}} > m_{H^\pm}$  then a new decay channel becomes available for  $H^{\pm\pm}$ , namely  $H^{\pm\pm} \rightarrow H^\pm W^*$ . Some attention has been given to the decay  $H^{\pm\pm} \rightarrow H^\pm W^*$ , and it has been shown that it can be the dominant channel over a wide range of values of  $\Delta M$  and  $v_\Delta$  [19, 23, 27, 33], even for  $\Delta M \ll m_W$ .

Another scenario is the case of  $m_{H^\pm} > m_{H^{\pm\pm}}$ , which would give rise to a new decay channel for the singly charged scalar, namely  $H^\pm \rightarrow H^{\pm\pm}W^*$ . This possibility has been mentioned in the context of the HTM in [23] only. We will perform the first study of the

magnitude of its branching ratio, as well as quantify its contribution to the production of  $H^{\pm\pm}$  at the LHC.<sup>1</sup> The decay rate for  $H^\pm \rightarrow H^{\pm\pm}W^*$  is easily obtained from the corresponding expression for the decay rate for  $H^{\pm\pm} \rightarrow H^\pm W^*$ , and thus one expects that  $H^\pm \rightarrow H^{\pm\pm}W^*$  will be sizeable over a wide range of values of  $\Delta M$  and  $v_\Delta$ . We point out for the first time that the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  would give rise to an alternative way to produce  $H^{\pm\pm}$  in pairs ( $H^{++}H^{--}$ ), namely by the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  followed by  $H^\mp \rightarrow H^{\pm\pm}W^*$ . Production of  $H^{++}H^{--}$  can give rise to a distinctive signature of four leptons ( $\ell^+\ell^+\ell^-\ell^-$ ), and simulations and searches of this channel currently only assume production via the process  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ .

Our work is organised as follows. In section II we describe the theoretical structure of the HTM. In section III the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  is introduced. Section IV contains our numerical analysis of the magnitude of the cross section for  $H^{++}H^{--}$  which originates from production via  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  followed by the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$ . Conclusions are given in section V.

## II. THE HIGGS TRIPLET MODEL

In the HTM [3, 6, 7] a  $Y = 2$  complex  $SU(2)_L$  isospin triplet of scalar fields is added to the SM Lagrangian. Such a model can provide Majorana masses for the observed neutrinos without the introduction of  $SU(2)_L$  singlet neutrinos via the gauge invariant Yukawa interaction:

$$\mathcal{L} = h_{\ell\ell'} L_\ell^T C i\tau_2 \Delta L_{\ell'} + \text{h.c.} \quad (1)$$

Here  $h_{\ell\ell'}(\ell, \ell' = e, \mu, \tau)$  is a complex and symmetric coupling,  $C$  is the Dirac charge conjugation operator,  $\tau_i$  is the Pauli matrix,  $L_\ell = (\nu_{\ell L}, \ell_L)^T$  is a left-handed lepton doublet, and  $\Delta$  is a  $2 \times 2$  representation of the  $Y = 2$  complex triplet fields:

$$\Delta = \begin{pmatrix} \Delta^+/\sqrt{2} & \Delta^{++} \\ \Delta^0 & -\Delta^+/\sqrt{2} \end{pmatrix}. \quad (2)$$

A non-zero triplet vacuum expectation value  $\langle \Delta^0 \rangle$  gives rise to the following mass matrix for neutrinos:

$$m_{\ell\ell'} = 2h_{\ell\ell'} \langle \Delta^0 \rangle = \sqrt{2}h_{\ell\ell'} v_\Delta. \quad (3)$$

The necessary non-zero  $v_\Delta$  arises from the minimisation of the most general  $SU(2)_L \otimes U(1)_Y$  invariant Higgs potential [7, 35], which is written<sup>2</sup> as follows [22, 23] (with  $\Phi = (\phi^+, \phi^0)^T$ ):

$$V = m^2(\Phi^\dagger \Phi) + \lambda_1(\Phi^\dagger \Phi)^2 + M^2 \text{Tr}(\Delta^\dagger \Delta) + \lambda_2[\text{Tr}(\Delta^\dagger \Delta)]^2 + \lambda_3 \text{Det}(\Delta^\dagger \Delta) \\ + \lambda_4(\Phi^\dagger \Phi) \text{Tr}(\Delta^\dagger \Delta) + \lambda_5(\Phi^\dagger \tau_i \Phi) \text{Tr}(\Delta^\dagger \tau_i \Delta) + \left( \frac{1}{\sqrt{2}} \mu (\Phi^T i\tau_2 \Delta^\dagger \Phi) + \text{h.c.} \right). \quad (4)$$

<sup>1</sup> The decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  has also been briefly mentioned in [34] in the context of a model with an isospin 3/2 multiplet with hypercharge  $Y = 3$ , which also includes triply charged Higgs bosons.

<sup>2</sup> One may rewrite the potential in eq. (4) by using  $2\text{Det}(\Delta^\dagger \Delta) = [\text{Tr}(\Delta^\dagger \Delta)]^2 - \text{Tr}[(\Delta^\dagger \Delta)^2]$  and  $(\Phi^\dagger \tau_i \Phi) \text{Tr}(\Delta^\dagger \tau_i \Delta) = 2\Phi^\dagger \Delta \Delta^\dagger \Phi - (\Phi^\dagger \Phi) \text{Tr}(\Delta^\dagger \Delta)$ .

Here  $m^2 < 0$  in order to ensure  $\langle \phi^0 \rangle = v/\sqrt{2}$  which spontaneously breaks  $SU(2) \otimes U(1)_Y$  to  $U(1)_Q$ , and  $M^2 (> 0)$  is the mass term for the triplet scalars. In the model of Gelmini-Roncadelli [35] the term  $\mu(\Phi^T i\tau_2 \Delta^\dagger \Phi)$  is absent, which leads to spontaneous violation of lepton number for  $M^2 < 0$ . The resulting Higgs spectrum contains a massless triplet scalar (majoron,  $J$ ) and another light scalar ( $H^0$ ). Pair production via  $e^+e^- \rightarrow H^0 J$  would give a large contribution to the invisible width of the  $Z$  and this model was excluded at the CERN Large Electron Positron Collider (LEP). The inclusion of the term  $\mu(\Phi^T i\tau_2 \Delta^\dagger \Phi)$  [7] explicitly breaks lepton number  $L\#$  when  $\Delta$  is assigned  $L\# = -2$ , and eliminates the majoron. Thus the scalar potential in eq. (4) together with the triplet Yukawa interaction of eq. (1) lead to a phenomenologically viable model of neutrino mass generation. For small  $v_\Delta/v$ , the expression for  $v_\Delta$  resulting from the minimisation of  $V$  is:

$$v_\Delta \simeq \frac{\mu v^2}{2M^2 + (\lambda_4 + \lambda_5)v^2} . \quad (5)$$

For large  $M$  compared to  $v$  one has  $v_\Delta \simeq \mu v^2/2M^2$ , which is sometimes referred to as the ‘‘Type II seesaw mechanism’’ and would naturally lead to a small  $v_\Delta$ . Recently there has been much interest in the scenario of light triplet scalars ( $M \approx v$ ) within the discovery reach of the LHC, for which eq. (5) leads to  $v_\Delta \approx \mu$ . In extensions of the HTM the term  $\mu(\Phi^T i\tau_2 \Delta^\dagger \Phi)$  may arise in various ways: i) it can be generated at tree level via the vev of a Higgs singlet field [36]; ii) it can arise at higher orders in perturbation theory [23]; iii) it can originate in the context of extra dimensions [22].

An upper limit on  $v_\Delta$  can be obtained from considering its effect on the parameter  $\rho (= M_W^2/M_Z^2 \cos^2 \theta_W)$ . In the SM  $\rho = 1$  at tree-level, while in the HTM one has (where  $x = v_\Delta/v$ ):

$$\rho \equiv 1 + \delta\rho = \frac{1 + 2x^2}{1 + 4x^2} . \quad (6)$$

The measurement  $\rho \approx 1$  leads to the bound  $v_\Delta/v \lesssim 0.03$ , or  $v_\Delta \lesssim 8 \text{ GeV}$ . Production mechanisms which depend on  $v_\Delta$  (i.e.  $pp \rightarrow W^{\pm*} \rightarrow W^\mp H^{\pm\pm}$  and fusion via  $W^{\pm*}W^{\pm*} \rightarrow H^{\pm\pm}$  [17, 37, 38]) are not competitive with the processes  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  and  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  at the energies of the Fermilab Tevatron, but such mechanisms can be the dominant source of  $H^{\pm\pm}$  at the LHC if  $v_\Delta = \mathcal{O}(1) \text{ GeV}$  and  $m_{H^{\pm\pm}} > 500 \text{ GeV}$ . At the 1-loop level,  $v_\Delta$  must be renormalised and explicit analyses lead to bounds on its magnitude similar to the above bound from the tree-level analysis, e.g. see [39].

The scalar eigenstates in the HTM are as follows: i) the charged scalars  $H^{\pm\pm}$  and  $H^\pm$ ; ii) the CP-even neutral scalars  $h^0$  and  $H^0$ ; iii) a CP-odd neutral scalar  $A^0$ . The doubly charged  $H^{\pm\pm}$  is entirely composed of the triplet scalar field  $\Delta^{\pm\pm}$ , while the remaining eigenstates are in general mixtures of the doublet and triplet fields. However, such mixing is proportional to the triplet vev, and hence small *even if*  $v_\Delta$  assumes its largest value of a few GeV.<sup>3</sup> Therefore  $H^\pm, H^0, A^0$  are predominantly composed of the triplet fields, while  $h^0$  is predominantly composed of the doublet field and plays the role of the SM Higgs boson. The scale of squared masses of  $H^{\pm\pm}, H^\pm, H^0, A^0$  are determined by  $M^2 + \lambda_4 v^2/2$  with mass splittings of

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<sup>3</sup> A large mixing angle is possible in the CP-even sector provided that  $m_{h^0} \sim m_{H^0}$  [40, 41].

order  $\lambda_5 v^2$  [22, 23, 41]:

$$\begin{aligned} m_{H^{\pm\pm}}^2 &\simeq m_{H^\pm}^2 - \frac{\lambda_5}{2} v^2, \\ m_{H^\pm}^2 &\simeq M^2 + \frac{\lambda_4}{2} v^2, \\ m_{H^0, A^0}^2 &\simeq m_{H^\pm}^2 + \frac{\lambda_5}{2} v^2. \end{aligned} \tag{7}$$

The degeneracy  $m_{H^0} \simeq m_{A^0}$  can be understood by the fact that the Higgs potential is invariant under a global  $U(1)$  for  $\Delta$  ( $L\#$  conservation) when one neglects the trilinear term proportional to  $\mu$ .

The mass hierarchy  $m_{H^{\pm\pm}} < m_{H^\pm} < m_{H^0, A^0}$  is obtained for  $\lambda_5 > 0$ , and the opposite hierarchy  $m_{H^{\pm\pm}} > m_{H^\pm} > m_{H^0, A^0}$  is obtained for  $\lambda_5 < 0$ . In general, one would not expect degenerate masses for  $H^{\pm\pm}, H^\pm, H^0, A^0$ , but instead one of the above two mass hierarchies. The sign of  $\lambda_5$  is not fixed by theoretical requirements of vacuum stability of the scalar potential [42], although  $|\lambda_5| < 2m_{H^\pm}^2/v^2$  is necessary to ensure that  $m_{H^{\pm\pm}}^2$  and  $m_{H^0, A^0}^2$  in eq. (8) are positive. Therefore the decays channels  $H^\pm \rightarrow H^{\pm\pm}W^*$  and  $H^{\pm\pm} \rightarrow H^\pm W^*$  are possible in the HTM.

### III. THE DECAY $H^\pm \rightarrow H^{\pm\pm}W^*$ AND PRODUCTION OF $H^{++}H^{--}$ FROM $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$

The potential importance of the decay channel  $H^\pm \rightarrow H^{\pm\pm}W^*$  (for  $m_{H^\pm} > m_{H^{\pm\pm}}$ ) has not been quantified in the HTM. For this decay to be kinematically open [23] one needs the mass hierarchy where  $H^{\pm\pm}$  is the lightest of the triplet scalars ( $m_{H^{\pm\pm}} < m_{H^\pm} < m_{H^0, A^0}$ ), which is obtained for  $\lambda_5 > 0$ . For the opposite mass hierarchy with  $\lambda_5 < 0$  ( $m_{H^{\pm\pm}} > m_{H^\pm} > m_{H^0, A^0}$ ) the related decay  $H^{\pm\pm} \rightarrow H^\pm W^*$  was shown to be important in the HTM in [19, 23, 27, 33]. The expression for the decay width of  $H^\pm \rightarrow H^{\pm\pm}W^*$  is easily obtained from the expression for  $H^{\pm\pm} \rightarrow H^\pm W^*$  by merely interchanging  $m_{H^{\pm\pm}}$  and  $m_{H^\pm}$ . After summing over all fermion states for  $W^* \rightarrow f'\bar{f}$ , excluding the  $t$  quark, the decay rate is given by

$$\Gamma(H^\pm \rightarrow H^{\pm\pm}W^* \rightarrow H^{\pm\pm}f'\bar{f}) \simeq \frac{9G_F^2 m_W^4 m_{H^\pm}}{4\pi^3} \int_0^{1-\kappa_{H^{\pm\pm}}} dx_2 \int_{1-x_2-\kappa_{H^{\pm\pm}}}^{1-\frac{\kappa_{H^{\pm\pm}}}{1-x_2}} dx_1 F_{H^{\pm\pm}W}(x_1, x_2), \tag{8}$$

where  $\kappa_{H^{\pm\pm}} \equiv m_{H^{\pm\pm}}/m_{H^\pm}$  and the analytical expression for  $F_{ij}(x_1, x_2)$  can be found in [43] (see also [44]). Note that this decay mode does not depend on  $v_\Delta$ . In eq. (8) we take  $f'$  and  $\bar{f}$  to be massless, which is a good approximation as long as the mass splitting between  $m_{H^{\pm\pm}}$  and  $m_{H^\pm}$  is above the mass of the charmed hadrons ( $\sim 2$  GeV). In our numerical analysis we will be mostly concerned with sizeable mass splittings,  $m_{H^\pm} - m_{H^{\pm\pm}} \gg 2$  GeV.

The other possible decays for  $H^\pm$  are  $H^\pm \rightarrow \ell^\pm \nu_{\ell'}$ ,  $H^\pm \rightarrow W^\pm Z$ ,  $H^\pm \rightarrow W^\pm h^0$  (where  $h^0$  is the SM-like scalar field) and  $H^\pm \rightarrow \bar{t}b$ . Explicit expressions for the decay widths of these channels can be found in the literature (e.g. [27, 37, 45]) and they are presented below. The decay width for  $H^\pm \rightarrow \ell^\pm \nu_{\ell'}$  is given by

$$\Gamma(H^\pm \rightarrow \ell^\pm \nu) \equiv \sum_{\ell, \ell'} \Gamma(H^\pm \rightarrow \ell^\pm \nu_{\ell'}) \simeq \frac{m_{H^\pm} \sum_i m_i^2}{16\pi v_\Delta^2}. \tag{9}$$

Note that  $\Gamma(H^\pm \rightarrow \ell^\pm \nu)$  has no dependence on the neutrino mixing angles because  $\sum_{\ell, \ell'} |h_{\ell \ell'}|^2 = \sum_i m_i^2 / (2v_\Delta^2)$ , where  $m_i$  ( $i = 1-3$ ) are neutrino masses. The decay widths for the channels which are proportional to  $v_\Delta^2$  are expressed as follows:

$$\Gamma(H^\pm \rightarrow W^\pm Z) \simeq \frac{v_\Delta^2 G_F^2 m_{H^\pm}^3}{4\pi} [\beta(m_{H^\pm}, m_W, m_Z)]^3, \quad (10)$$

$$\Gamma(H^\pm \rightarrow W^\pm h^0) \simeq \frac{v_\Delta^2 G_F^2 m_{H^\pm}^3}{4\pi} \left( \frac{2m_{H^\pm}^2 - \lambda_4 v^2}{m_{H^\pm}^2 - m_h^2} - 1 \right)^2 [\beta(m_{H^\pm}, m_W, m_h)]^3, \quad (11)$$

$$\Gamma(H^\pm \rightarrow \bar{t}b) \simeq \frac{3v_\Delta^2 G_F^2 m_t^2 m_{H^\pm}}{2\pi} \left( 1 - \frac{m_t^2}{m_{H^\pm}^2} \right)^2, \quad (12)$$

$$\beta(m_1, m_2, m_3) \equiv \sqrt{1 - \frac{(m_2 + m_3)^2}{m_1^2}} \sqrt{1 - \frac{(m_2 - m_3)^2}{m_1^2}}. \quad (13)$$

The decay  $H^\pm \rightarrow W^\pm h^0$  is caused by two small mixings of scalar fields. One is the mixing angle  $\theta_\pm \simeq \sqrt{2}v_\Delta/v$  between  $\phi^\pm$  and  $\Delta^\pm$ , and the other is the mixing angle  $\theta_0 \simeq (2m_{H^\pm}^2 - \lambda_4 v^2)v_\Delta / ((m_{H^0}^2 - m_h^2)v)$  between  $\text{Re}(\phi^0)$  and  $\text{Re}(\Delta^0)$ . If  $M \gg v$ , then one has  $(2m_{H^\pm}^2 - \lambda_4 v^2)/(m_{H^0}^2 - m_h^2) \simeq 2$  in eq. (11). Since we are interested in the case where the exotic scalars have masses of the electroweak scale, we do not take a very large  $M$ . However, we assume  $(2m_{H^\pm}^2 - \lambda_4 v^2)/(m_{H^0}^2 - m_h^2) \simeq 2$  for simplicity, which can be achieved by  $\lambda_4 \simeq 2(m_{H^\pm}^2 - m_{H^0}^2 + m_h^2)/v^2$ . The decay  $H^\pm \rightarrow \bar{t}b$  is mediated by the small  $\phi^\pm$  component of  $H^\pm$  through  $\theta_\pm$ . For  $m_{H^\pm} = \mathcal{O}(100)$  GeV,  $\Gamma(H^\pm \rightarrow \bar{t}b) \propto m_t^2 m_{H^\pm}$  is comparable to  $\Gamma(H^\pm \rightarrow W^\pm Z)$  and  $\Gamma(H^\pm \rightarrow W^\pm h^0) \propto m_{H^\pm}^3$ . These three decay widths in eq. (10)-(12) are greater than  $\Gamma(H^\pm \rightarrow \ell \nu)$  for  $v_\Delta \gtrsim 0.1$  MeV while  $\Gamma(H^\pm \rightarrow \ell \nu)$  dominates for  $v_\Delta \lesssim 0.1$  MeV.

It has already been shown that the decay  $H^{\pm\pm} \rightarrow H^\pm W^*$  can be the dominant decay channel for the doubly charged scalar over a wide range of values of  $\Delta M \equiv m_{H^{\pm\pm}} - m_{H^\pm}$  and  $v_\Delta$  [19, 23, 27, 33], even for  $\Delta M \ll m_W$ . Hence we expect a similar result for the decay  $H^\pm \rightarrow H^{\pm\pm} W^*$  for the singly charged scalar. The branching ratio  $\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)$  will be maximised with respect to  $v_\Delta$  if  $\Gamma(H^\pm \rightarrow \ell^\pm \nu) = \Gamma(H^\pm \rightarrow W^\pm Z) + \Gamma(H^\pm \rightarrow W^\pm h^0) + \Gamma(H^\pm \rightarrow \bar{t}b)$  which is achieved for  $v_\Delta \simeq 0.1$  MeV. A numerical study of the magnitude of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)$  is presented in the next section.

We now emphasise an important phenomenological difference between the distinct scenarios of a sizeable branching ratio for the decay channels  $H^{\pm\pm} \rightarrow H^\pm W^*$  (for  $\lambda_5 < 0$ ) and  $H^\pm \rightarrow H^{\pm\pm} W^*$  (for  $\lambda_5 > 0$ ). The decay  $H^{\pm\pm} \rightarrow H^\pm W^*$  is expected to weaken the discovery potential of  $H^{\pm\pm}$  at the LHC, because it would reduce the branching ratio of a channel like  $H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm$  (which is otherwise the dominant channel for  $v_\Delta \lesssim 0.1$  MeV, and enjoys low SM backgrounds). We note that there has been no simulation of the detection prospects of  $H^{\pm\pm} \rightarrow H^\pm W^*$ , and its signature would be different to that of the standard decay channels  $H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm$  and  $H^{\pm\pm} \rightarrow W^\pm W^\pm$ .

In contrast, we point out that the decay  $H^\pm \rightarrow H^{\pm\pm} W^*$  could actually *improve* the discovery potential of  $H^{\pm\pm}$  at the LHC. From the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm} H^\mp$  the decay mode  $H^\pm \rightarrow H^{\pm\pm} W^*$  would give rise to pair production ( $H^{++}H^{--}$ ) of doubly charged Higgs bosons. We believe that this additional way to produce  $H^{\pm\pm}$  has not been



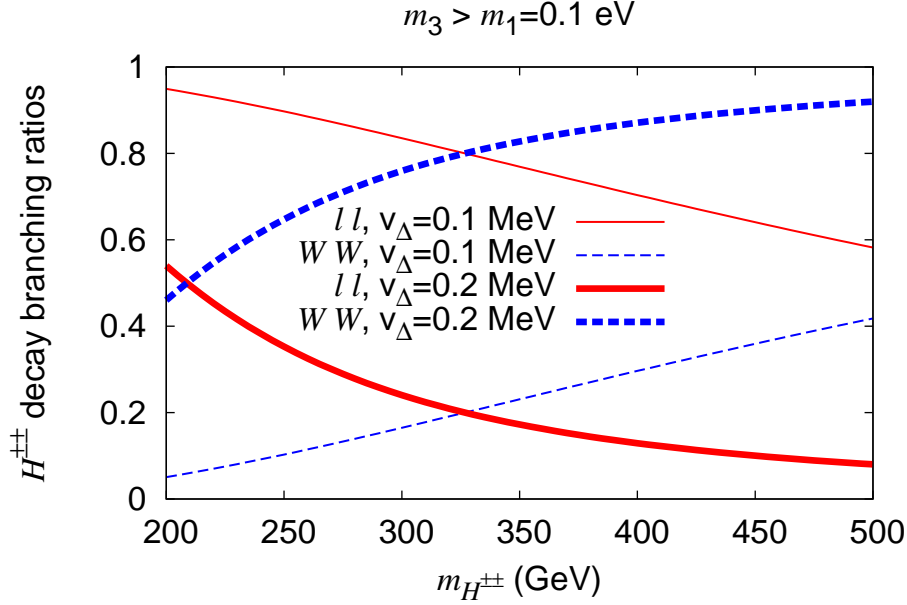


FIG. 1:  $\sum_{\ell, \ell'} \text{BR}(H^{\pm\pm} \rightarrow \ell^{\pm} \ell'^{\pm})$  with red solid lines and  $\text{BR}(H^{\pm\pm} \rightarrow W^{\pm} W^{\pm})$  with blue dashed lines as a function of  $m_{H^{\pm\pm}} (\leq m_{H^{\pm}})$  for  $v_{\Delta} = 0.1$  MeV (thin lines) and  $v_{\Delta} = 0.2$  MeV (bold lines). For the neutrino masses we used  $m_1 = 0.1$  eV with  $\Delta m_{31}^2 > 0$ .

discussed before. In this scenario  $H^{\pm\pm}$  is the lightest of the triplet scalars, and its only possible decay channels are  $H^{\pm\pm} \rightarrow \ell^{\pm} \ell'^{\pm}$  and  $H^{\pm\pm} \rightarrow W^{\pm} W^{\pm}$ , with branching ratios determined by the magnitude of  $v_{\Delta}$ . These two branching ratios can be of the same order of magnitude for  $v_{\Delta} \simeq 0.1$  MeV, as can be seen in Fig. 1 where we fix  $v_{\Delta} = 0.1$  MeV and 0.2 MeV (similar figures can be found in [27]). In the range of  $m_{H^{\pm\pm}} = 200 \rightarrow 500$  GeV, one has  $\sum_{\ell, \ell'} \text{BR}(H^{\pm\pm} \rightarrow \ell^{\pm} \ell'^{\pm}) \simeq 100\%$  for  $v_{\Delta} \lesssim 0.05$  MeV, while for  $v_{\Delta} \gtrsim 0.4$  MeV one has  $\text{BR}(H^{\pm\pm} \rightarrow W^{\pm} W^{\pm}) \simeq 100\%$ .

In simulations of pair production of  $H^{\pm\pm}$  it is assumed that the production channel  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++} H^{--}$  is the only mechanism. If  $v_{\Delta} \lesssim 0.1$  MeV then the decay channel  $H^{\pm\pm} \rightarrow \ell^{\pm} \ell'^{\pm}$  is dominant, and four-lepton signatures ( $4\ell$ ) would be possible. Studies have shown that the Standard Model background for the  $4\ell$  signature [29] is considerably smaller than that for the signature of  $3\ell$  [30, 31], and at present it is assumed that the  $4\ell$  signature can only arise from  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++} H^{--}$ . The importance of the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm} H^{\mp}$  has been appreciated for the  $3\ell$  signature, in which the decay  $H^{\mp} \rightarrow \ell^{\mp} \nu$  is assumed [19, 27, 30, 31, 46]. For the case of a sizeable branching ratio for  $H^{\pm} \rightarrow H^{\pm\pm} W^*$  we point out that the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm} H^{\mp}$  can also contribute to the  $4\ell$  signature, which is the signature with lowest background. Searches for four leptons originating from  $H^{++} H^{--}$  have already been performed by the Tevatron [21] and LHC [32]. If  $\text{BR}(H^{\pm} \rightarrow H^{\pm\pm} W^*)$  were sizeable we would expect a strengthening of the derived limit on  $m_{H^{\pm\pm}}$ .

## IV. NUMERICAL ANALYSIS

In this section we quantify the magnitude of the number of pair-produced  $H^{++}H^{--}$  arising from the process  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  with decay  $H^\pm \rightarrow H^{\pm\pm}W^*$ , and make a comparison with the conventional mechanism  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ .

The important parameters for our analyses are  $v_\Delta$ ,  $m_{H^\pm}$ , and  $m_{H^{\pm\pm}}$ . We take  $m_{H^{\pm\pm}} = 200$  GeV or 500 GeV and show results as functions of  $v_\Delta$  and  $m_{H^\pm}$ . The decay branching ratios of  $H^\pm$  also depend on two undetermined parameters,  $m_h$  and  $m_1$  (one of the neutrino masses). These are fixed as  $m_h = 120$  GeV and  $m_1 = 0.1$  eV in our numerical analysis. Note that  $m_h$  only enters through the decay width for  $H^\pm \rightarrow W^\pm h^0$ . Neutrino oscillation experiments [8–12] provide a measurement of two neutrino mass differences,  $\Delta m_{ij}^2 \equiv m_i^2 - m_j^2$ , and we use the following values:  $\Delta m_{21}^2 = 7.6 \times 10^{-5}$  eV,  $|\Delta m_{31}^2| = 2.4 \times 10^{-3}$  eV. Although  $\Delta m_{31}^2 > 0$  (referred to as “normal mass ordering”) is also assumed in our analysis, our results do not change significantly for  $\Delta m_{31}^2 < 0$  because the neutrino masses are almost degenerate for  $m_1 = 0.1$  eV.

The experimental bound  $\text{BR}(\mu \rightarrow \bar{e}ee) < 1.0 \times 10^{-12}$  gives a stringent constraint on  $|h_{e\ell}|$  and  $m_{H^{\pm\pm}}$ .<sup>4</sup> Assuming naively  $|m_{ee}| \simeq |m_{e\mu}| \simeq m_1$  for  $m_1 = 0.1$  eV, the bound on  $\text{BR}(\mu \rightarrow \bar{e}ee) = |m_{ee}|^2 |m_{e\mu}|^2 / (16G_F^2 v_\Delta^4 m_{H^{\pm\pm}}^4)$  can be translated into the constraint  $v_\Delta m_{H^{\pm\pm}} \gtrsim 1.5 \times 10^4$  eV · GeV. Therefore, we use  $v_\Delta \geq 100$  eV in order to satisfy this constraint for  $m_{H^{\pm\pm}} = 200$  GeV.

In Fig. 2 we show the BRs of  $H^\pm$  decays into  $H^{\pm\pm}W^*$  (red solid),  $\ell\nu$  (blue dashed),  $\bar{t}b$  (green dotted),  $WZ$  (magenta dot-dashed), and  $Wh^0$  (cyan dot-dot-dashed) as a function of  $m_{H^\pm}$  for various values of  $v_\Delta$ , fixing  $m_{H^{\pm\pm}} = 200$  GeV and  $m_{h^0} = 120$  GeV. The range of  $m_{H^\pm}$  in the figures corresponds to  $0 \leq \lambda_5 \lesssim 1$ , which easily satisfies the perturbative constraint  $\lambda_5 < 4\pi$ . Very large mass splittings (e.g.  $\gg 100$  GeV) are constrained by measurements of electroweak precision observables, but the mass splittings in Fig. 2 are compatible with the analyses in [39] (which are for models with a  $Y = 0$  triplet). In Fig. 2(a) we fix  $v_\Delta = 100$  eV, for which  $\sum_{\ell, \ell'} \text{BR}(H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm) \simeq 100\%$ . One can see that  $H^\pm \rightarrow H^{\pm\pm}W^*$  competes with  $H^\pm \rightarrow \ell^\pm \nu$ , with all other decay channels being negligible. For  $|\Delta M| > 20$  GeV,  $H^\pm \rightarrow H^{\pm\pm}W^*$  becomes the dominant decay channel. In Fig. 2(b) we fix  $v_\Delta = 0.1$  MeV, and  $H^\pm \rightarrow H^{\pm\pm}W^*$  becomes the dominant decay channel for much smaller mass splittings,  $|\Delta M| > 2$  GeV. In Fig. 2(c) we fix  $v_\Delta = 1$  GeV, for which the competing decays are  $H^\pm \rightarrow tb$ ,  $H^\pm \rightarrow WZ$  and  $H^\pm \rightarrow Wh^0$ . In this scenario the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  becomes the dominant channel for  $|\Delta M| > 30$  GeV.

In Fig. 3 we show contours of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*)$  in the plane  $[m_{H^\pm}, v_\Delta]$ . The red solid, green dashed, and blue dotted lines correspond to contours of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*) = 0.5, 0.9, \text{ and } 0.99$ , respectively. The BR is maximised at around  $v_\Delta = 0.1$  MeV, as expected. It is clear from Fig. 2 and Fig. 3 that the decay of  $H^\pm$  into  $H^{\pm\pm}$  can be dominant in a wide region of the parameter space of the HTM even if the two-body decay into  $H^{\pm\pm}W^\mp$  (for  $m_{H^\pm} > m_{H^{\pm\pm}} + m_W$ ) is forbidden kinematically. Moreover, for  $v_\Delta = 100$  eV (i.e. when the four-lepton signal arising from the decay of  $H^{++}H^{--}$  is dominant) the magnitude of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*)$  becomes very large if  $|\Delta M| \gtrsim 30$  GeV.

We now study the magnitude of the number of pair-produced  $H^{++}H^{--}$  which originate from  $pp \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  followed by the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$ . We define the variable

<sup>4</sup> This stringent constraint can be avoided if  $|m_{e\mu}| = 0$  [23] or  $|m_{ee}| = 0$  [47]. See also [48].



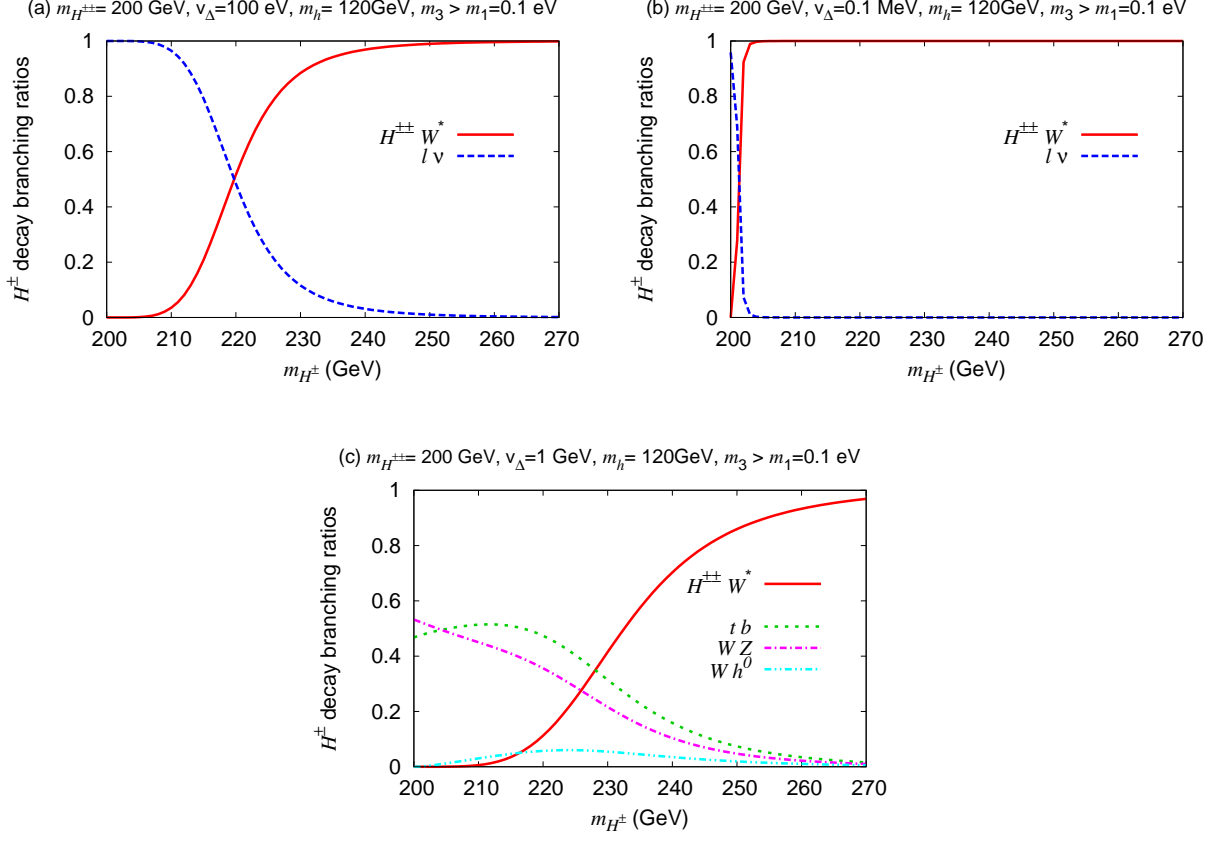


FIG. 2: The BRs of  $H^\pm$  decays into  $H^{\pm\pm}W^*$  (red solid line),  $\ell\nu$  (blue dashed line),  $t\bar{b}$  (green dotted line),  $WZ$  (magenta dot-dashed line), and  $Wh^0$  (cyan dot-dot-dashed line) as a function of  $m_{H^\pm}$ . In all panels  $m_{H^{\pm\pm}} = 200$  GeV and  $m_{h^0} = 120$  GeV. In panels a), b) and c) we fix  $v_\Delta = 100$  eV, 0.1 MeV and 1 GeV respectively. For the neutrino masses we used  $m_1 = 0.1$  eV with  $\Delta m_{31}^2 > 0$ . In each panel the decay modes with a negligible BR are omitted.

$X_1$  as follows:

$$X_1 \equiv \left\{ \sigma(pp \rightarrow W^* \rightarrow H^{++}H^-) + \sigma(pp \rightarrow W^* \rightarrow H^{--}H^+) \right\} BR(H^\pm \rightarrow H^{\pm\pm}W^*). \quad (14)$$

In Fig. 4 we show the behaviour of  $\sigma(H^{++}H^{--}) \equiv \sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}) + X_1$  with respect to  $m_{H^\pm}$  for several values of  $v_\Delta$ . In Fig. 4a we take  $m_{H^{\pm\pm}} = 200$  GeV and  $\sqrt{s} = 7$  TeV, and in Fig. 4b we take  $m_{H^{\pm\pm}} = 500$  GeV and  $\sqrt{s} = 14$  TeV. We use CTEQ6L1 parton distribution functions [49]. The range of  $m_{H^\pm}$  in Fig. 4b corresponds to  $0 \leq \lambda_5 \lesssim 2.5$ . The horizontal dot-dashed line corresponds to the case of  $X_1 = 0$ , i.e. the magnitude of  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--})$  alone. The red solid, green dashed, and blue dotted lines are the results with  $v_\Delta = 100$  eV, 0.1 MeV, and 1 GeV, respectively. The red solid line (for which  $\sum_{\ell, \ell'} BR(H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm) \simeq 100\%$ ) shows that the extra contribution from  $H^\pm \rightarrow H^{\pm\pm}W^*$  can enhance the number of four-lepton events by a factor of 2 (at  $m_{H^\pm} \simeq 230$  GeV in Fig. 4a) and 2.4 (at  $m_{H^\pm} \simeq 540$  GeV in Fig. 4b). For  $v_\Delta = 0.1$  MeV, around which  $\sum_{\ell, \ell'} BR(H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm)$  can still be sizeable (See Fig. 1), the enhancement factor for pair-produced  $H^{++}H^{--}$  can be as large as 2.6 in Fig. 4a and 2.8 in Fig. 4b. For  $v_\Delta = 1$  GeV

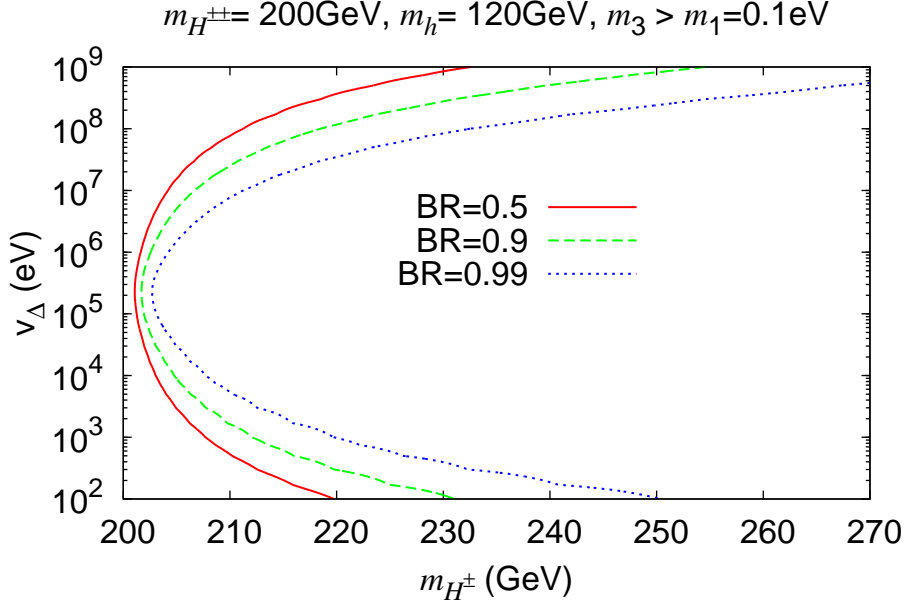


FIG. 3: Contours of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*)$  in the plane  $(m_{H^\pm}, v_\Delta)$ . We used  $m_1 = 0.1 \text{ eV}$  with  $\Delta m_{31}^2 > 0$ . The red solid, green dashed, and blue dotted lines show contours of  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*) = 0.5, 0.9$  and  $0.99$ , respectively.

the enhancement of pair-produced  $H^{++}H^{--}$  is interpreted as an increase in the number of  $W^+W^+W^-W^-$  events, because  $\text{BR}(H^{\pm\pm} \rightarrow W^\pm W^\pm) \simeq 100\%$ . The shape of the curves is caused by the different dependence of the cross section and BR on the mass splitting  $\Delta M$ . As  $m_{H^\pm}$  increases, the cross section of  $pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  is unaffected but the cross section of  $pp \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  decreases. However, a larger mass splitting is favourable from the point of view of the BR.

Finally, we note that a pair of  $H^{\pm\pm}$  can also be produced from other production mechanisms, namely  $H^+H^-$ ,  $H^\pm H^0$ ,  $H^\pm A^0$ , and  $H^0 A^0$ . Although the contribution from  $H^{\pm\pm}H^\mp$  in eq. (14) is the most important one because of the mass hierarchy  $m_{H^{\pm\pm}} < m_{H^\pm} < m_{H^0, A^0}$  and its linear dependence on  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*)$ , the above mechanisms can give a significant contribution to the number of pair-produced  $H^{\pm\pm}$ , as will be described qualitatively below.

Naively, one would expect the next most important mechanism to be  $H^+H^-$  because its contribution to the production of  $H^{++}H^{--}$  scales as  $\text{BR}^2$  as follows:

$$X_2 \equiv \sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^+H^-) [\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*)]^2. \quad (15)$$

However, the couplings for  $\gamma H^+H^-$  and  $Z H^+H^-$  are about a half of those for  $\gamma H^{++}H^{--}$  and  $Z H^{++}H^{--}$ , respectively. The interference between  $\gamma^*$  and  $Z^*$  is destructive for  $H^+H^-$  production while it is constructive for  $H^{++}H^{--}$  production. Tables I and II show that  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^+H^-)$  is smaller than  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--})$  by a factor of  $\sim 9.5$ , even for  $m_{H^\pm} = m_{H^{\pm\pm}}$  (see e.g. [50]). Moreover,  $X_2$  is suppressed relative to  $X_1$  by an extra factor of BR when  $\text{BR}(H^\pm \rightarrow H^{\pm\pm}W^*) \neq 100\%$ . Therefore the contribution from  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^+H^-)$  to the production of  $H^{++}H^{--}$  is considerably less than the QCD  $K$  factor for  $pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  (which is known to be around 1.25 at the LHC [15]).

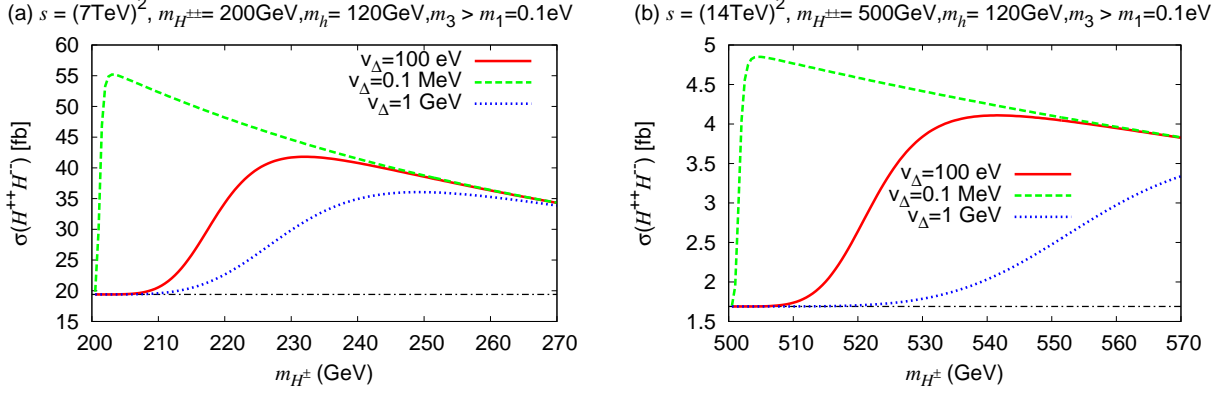


FIG. 4: Behaviour of  $\sigma(H^{++}H^{--}) \equiv \sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}) + X_1$  as a function of  $m_{H^\pm}$ . (a) for  $m_{H^\pm} = 200$  GeV at the LHC with  $\sqrt{s} = 7$  TeV. (b) for  $m_{H^\pm} = 500$  GeV at the LHC with  $\sqrt{s} = 14$  TeV. We used  $m_h = 120$  GeV and  $m_1 = 0.1$  eV with  $\Delta m_{31}^2 > 0$ . The red solid, green dashed, and blue dotted lines show results for  $v_\Delta = 100$  eV, 0.1 MeV, and 1 GeV, respectively. The horizontal dot-dashed line shows  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--})$ .

It turns out that the production of  $H^\pm H^0$  and  $H^\pm A^0$  are numerically more important than  $H^+ H^-$ , despite their contributions scaling as  $\text{BR}^3$ . The narrow width approximation for contributions from  $H^0$  and  $A^0$  with  $m_{H^0} \simeq m_{A^0}$  is rather complicated because of their interference. We define the variables  $X_3$  and  $X'_3$  as follows:

$$X_3 \equiv \{ \sigma(pp \rightarrow W^* \rightarrow H^+ H^0) + \sigma(pp \rightarrow W^* \rightarrow H^- H^0) \} \times \text{BR}_+ [\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)]^2, \quad (16)$$

$$X'_3 \equiv \{ \sigma(pp \rightarrow W^* \rightarrow H^+ H^0) + \sigma(pp \rightarrow W^* \rightarrow H^- H^0) \} \times \text{BR}_- [\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)]^2, \quad (17)$$

$$\begin{aligned} \text{BR}_\pm &\equiv \text{BR}(H^0 \rightarrow H^\pm W^*) + \text{BR}(A^0 \rightarrow H^\pm W^*) \\ &\pm \frac{4\text{BR}(H^0 \rightarrow H^\pm W^*)\text{BR}(A^0 \rightarrow H^\pm W^*)}{\text{BR}(H^0 \rightarrow H^\pm W^*) + \text{BR}(A^0 \rightarrow H^\pm W^*)}, \end{aligned} \quad (18)$$

where we used  $\sigma(pp \rightarrow W^* \rightarrow H^\pm A^0) \simeq \sigma(pp \rightarrow W^* \rightarrow H^\pm H^0)$  because  $m_{A^0} \simeq m_{H^0}$ . The interesting point is that  $X'_3$  is for the process which gives *same-sign*  $H^{++}H^{++}$  (with  $3(W^-)^*$ ) and  $H^{--}H^{--}$  (with  $3(W^+)^*$ ) while  $X_3$  is for  $H^{++}H^{--}$  production. Since  $X'_3$  arises as the breaking effect of the lepton number ( $\Delta$  has  $L\# = -2$ ), it vanishes for  $v_\Delta \rightarrow 0$ , for which the total decay widths satisfy  $\Gamma_{\text{tot}}(H^0) = \Gamma_{\text{tot}}(A^0)$ , namely  $\text{BR}(H^0 \rightarrow H^\pm W^*) = \text{BR}(A^0 \rightarrow H^\pm W^*)$ . This means that the *same-sign*  $H^{\pm\pm}H^{\pm\pm}$  would not give the *same-sign*  $4\ell$  signal because  $\text{BR}(H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm)$  is small for a large  $v_\Delta$  where  $X'_3$  could be sizeable. A pair of  $H^{\pm\pm}$  (*same-sign* or different sign) is provided by  $X_3 + X'_3$ , which is proportional to  $2[\text{BR}(H^0 \rightarrow H^\pm W^*) + \text{BR}(A^0 \rightarrow H^\pm W^*)]$ ; the factor of 2 compensates the fact that the sum of the cross sections in eq. (16) is a half of the sum in eq. (14) for  $m_{H^0, A^0} = m_{H^\pm} = m_{H^{\pm\pm}}$  as shown in Tables I and II. Although  $\text{BR}(A^0 \rightarrow H^+ W^*) = \text{BR}(A^0 \rightarrow H^- W^*)$  (likewise for  $H^0$ ) and the maximum value of each is 50%, this is compensated by  $\text{BR}(H^0 \rightarrow H^\pm W^*) + \text{BR}(A^0 \rightarrow H^\pm W^*)$  in  $X_3 + X'_3$ . Since the partial decay widths of  $H^0$  and  $A^0$  depend on the scalar masses and  $v_\Delta$  in a way which is very similar to the partial decay widths of  $H^\pm$  (see e.g. [27]), the

$\sqrt{s} = 7 \text{ TeV}$	$\sigma(pp \rightarrow V^* \rightarrow H_1 H_2) [\text{fb}]$						
$m_{H^{\pm\pm}} = 200 \text{ GeV}$	$H^{++}H^{--}$	$H^{++}H^-$	$H^{--}H^+$	$H^+H^-$	$H^+H^0$ (or $H^+A^0$ )	$H^-H^0$ (or $H^-A^0$ )	$H^0A^0$
$m_{H^\pm} = 200 \text{ GeV}$	19	26	10	2.0	13	5.1	18
$m_{H^\pm} = 230 \text{ GeV}$	19	19	7.3	1.1	5.6	2.1	6.0
$m_{H^\pm} = 270 \text{ GeV}$	19	13	4.8	0.54	2.2	0.76	1.9

TABLE I: Production cross sections of a pair of exotic Higgs bosons ( $H_1 H_2$ ) from off-shell gauge bosons ( $V^*$ ) in the HTM at the LHC with  $\sqrt{s} = 7 \text{ TeV}$ . We take  $m_{H^{\pm\pm}} = 200 \text{ GeV}$  and we use a relation  $m_{H^0, A^0}^2 = 2m_{H^\pm}^2 - m_{H^{\pm\pm}}^2$ ;  $m_{H^0, A^0} = 200, 257, 325 \text{ GeV}$  for  $m_{H^\pm} = 200, 230, 270 \text{ GeV}$ , respectively.

analogies of Fig. 3 for  $\text{BR}(A^0 \rightarrow H^\pm W^*)$  and  $\text{BR}(H^0 \rightarrow H^\pm W^*)$  would show a similar quantitative behaviour as Fig. 3.<sup>5</sup> Thus the main difference between  $X_1$  and  $X_3 + X'_3$  would be the phase space factor because we take  $m_{H^0, A^0} > m_{H^\pm}$ . The contribution of  $X_3 + X'_3$  to the production of a pair of  $H^{\pm\pm}$  would be sizeable for  $v_\Delta \simeq 0.1 \text{ MeV}$ , where the relevant BRs in eq. (16) could be very large for a small mass splitting. Moreover, the contribution of  $X_3 + X'_3$  would not be so small even for large mass splittings e.g.  $m_{H^\pm} = 270 \text{ GeV}$  and  $m_{H^{\pm\pm}} = 200 \text{ GeV}$  (which give  $m_{H^0, A^0} = 325 \text{ GeV}$ ), for which the BRs in eq. (16) could be maximal.

The last mechanisms (which scale as  $\text{BR}^4$ ) are

$$X_4 \equiv \sigma(pp \rightarrow Z^* \rightarrow H^0 A^0) \text{BR}_+^2 [\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)]^2, \quad (19)$$

$$X'_4 \equiv \sigma(pp \rightarrow Z^* \rightarrow H^0 A^0) \text{BR}_+ \text{BR}_- [\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)]^2, \quad (20)$$

$$X''_4 \equiv \sigma(pp \rightarrow Z^* \rightarrow H^0 A^0) \text{BR}_-^2 [\text{BR}(H^\pm \rightarrow H^{\pm\pm} W^*)]^2. \quad (21)$$

Note that  $X'_4$  gives a pair of *same-sign*  $H^{\pm\pm}$  (being proportional to  $\text{BR}_-$ , like  $X'_3$ ) and its magnitude is negligible for small  $v_\Delta$ . Although both of  $X_4$  and  $X''_4$  give  $H^{++}H^{--}$ ,  $X''_4$  also vanishes for  $v_\Delta \rightarrow 0$  because it is sensitive to  $\text{BR}_-^2$  i.e. it is quadratic in lepton number violation. The phase space suppression ( $m_{H^{\pm\pm}} < m_{H^\pm} < m_{H^0, A^0}$ ) ensures that  $\sigma(pp \rightarrow Z^* \rightarrow H^0 A^0)$  is much smaller than  $\sigma(pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--})$  for the case of a large mass splitting with  $\sqrt{s} = 7 \text{ TeV}$ . Therefore, for  $X_4$  to be important a large mass splitting with  $\sqrt{s} = 14 \text{ TeV}$  or a small mass splitting for  $v_\Delta \simeq 0.1 \text{ MeV}$  are preferred.

We note that the detection efficiencies for the above mechanisms ( $X_1, X_2, X_3$  and  $X_4$ ) would in general be different from that of the well-studied mechanism  $pp \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$  because of the extra  $W^*$ . We defer a detailed study to a future work.

<sup>5</sup> We note that the decays  $A^0 \rightarrow H^\pm W^*$  and  $H^0 \rightarrow H^\pm W^*$  were also mentioned as a source of  $H^\pm$  in [23].

$\sqrt{s} = 14 \text{ TeV}$	$\sigma(pp \rightarrow V^* \rightarrow H_1 H_2) [\text{fb}]$						
$m_{H^{\pm\pm}} = 500 \text{ GeV}$	$H^{++}H^{--}$	$H^{++}H^-$	$H^{--}H^+$	$H^+H^-$	$H^+H^0$ (or $H^+A^0$ )	$H^-H^0$ (or $H^-A^0$ )	$H^0A^0$
$m_{H^\pm} = 500 \text{ GeV}$	1.7	2.3	0.83	0.18	1.1	0.42	1.5
$m_{H^\pm} = 540 \text{ GeV}$	1.7	1.9	0.69	0.13	0.69	0.24	0.78
$m_{H^\pm} = 570 \text{ GeV}$	1.7	1.7	0.60	0.097	0.49	0.17	0.50

TABLE II: Production cross sections of a pair of exotic Higgs bosons ( $H_1 H_2$ ) from off-shell gauge bosons ( $V^*$ ) in the HTM at the LHC with  $\sqrt{s} = 14 \text{ TeV}$ . We take  $m_{H^{\pm\pm}} = 500 \text{ GeV}$  and we use a relation  $m_{H^0, A^0}^2 = 2m_{H^\pm}^2 - m_{H^{\pm\pm}}^2$ ;  $m_{H^0, A^0} = 500, 577, 632 \text{ GeV}$  for  $m_{H^\pm} = 500, 540, 570 \text{ GeV}$ , respectively.

## V. CONCLUSIONS

Doubly charged Higgs bosons ( $H^{\pm\pm}$ ), which arise in the Higgs Triplet Model (HTM) of neutrino mass generation, are being searched for at the Tevatron and at the LHC. We showed that  $H^{\pm\pm}$  can be produced from the decay of a singly charged Higgs boson ( $H^\pm$ ) via  $H^\pm \rightarrow H^{\pm\pm}W^*$ , which can have a large branching ratio in a wide region of the parameter space of the HTM. From the production mechanism  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$ , the above decay would give rise to pair production  $H^{++}H^{--}$ , with a number of events which can be comparable to that from the conventional mechanism  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ . Current simulations and searches for  $H^{++}H^{--}$  at the Tevatron/LHC assume production solely from  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ . The contribution from  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  with decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  would be an additional source of pair-produced  $H^{\pm\pm}$ , which should enhance the detection prospects in this channel (e.g. four-lepton signatures if the decay mode  $H^{\pm\pm} \rightarrow \ell^\pm \ell'^\pm$  is dominant). We also pointed out that production mechanisms involving the neutral triplet scalars ( $H^0, A^0$ ) of the HTM can contribute to pair production  $H^{++}H^{--}$  through the decay chain  $H^0, A^0 \rightarrow H^\pm W^*$  followed by  $H^\pm \rightarrow H^{\pm\pm}W^*$ . We advocate dedicated simulations of  $q'\bar{q} \rightarrow W^* \rightarrow H^{\pm\pm}H^\mp$  with the decay  $H^\pm \rightarrow H^{\pm\pm}W^*$  (and the analogous mechanisms with neutral scalars), and a comparison with  $q\bar{q} \rightarrow \gamma^*, Z^* \rightarrow H^{++}H^{--}$ .

## Acknowledgements

We thank Mayumi Aoki and Koji Tsumura for useful discussions. A.G.A was supported by a Marie Curie Incoming International Fellowship, FP7-PEOPLE-2009-IIF, Contract No. 252263. The work of H.S. was supported in part by the Sasakawa Scientific Research Grant from the Japan Science Society and Grant-in-Aid for Young Scientists (B) No. 23740210.

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